

# Toward the non-perturbative description of high energy processes

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## Abstract

General implications of existence of non-perturbative scales and hadronic sub-structure for high energy processes are discussed. We propose that the dependence of the cross section of  $\bar{q}q$  dipoles on their size  $d$  should deviate from  $d^2$  when  $d$  becomes comparable to substructure scale. Then we discuss Kharzeev-Levin pomeron model [1], based on ladder-type diagrams with *scalar* resonances (scalar  $\pi\pi$  or  $\sigma$  and the scalar glueball  $G_0$ ). This channel is truly unique, because instanton-induced attractive  $gg$  interaction [2] leads to unusually small sizes and strong coupling constants of these states, supplemented by unusually large mass scale,  $M_0 \approx 4 GeV$ , of the transition boundary to the perturbative regime. As pomeron is a small-size object by itself, these resonances may play a special role in its dynamics. Furthermore, we use more realistic description of the scalar gluonic spectral density without free parameters, and slightly modify the model to get correct chiral limit. We conclude that the non-perturbative part of the scalar contribution to the soft pomeron intercept is  $\Delta = .05 \pm 0.015$ , with comparable contributions from both  $\sigma$  and  $G_0$ .

1. Significant progress of non-perturbative QCD has been mostly related with approaches based on its Euclidean formulations: lattice simulations, semi-classical theory based on instantons etc. During the last decade we have learned a lot about correlation functions and their spectral densities, hadronic wave functions and form-factors. Dramatically different features of different channels pointed out in [3] was confirmed and studied in details [5]: below we discuss one of the most striking cases, the gluonic  $J^{PC} = O^{++}$  channel.

However, little of this progress has so far contributed toward understanding of high energy processes. True, it is difficult to translate many of those tools into Minkowski space. Elsewhere [25] we will report some semi-classical calculations aiming to bridge this gap, while we start this Letter with more general discussion of some qualitative ideas related hadronic substructure to high energy scattering.

The most important lesson we learned is that the non-perturbative objects in the QCD vacuum (and inside hadrons) are not some shapeless soft fields, with typical momenta of the order of  $\Lambda_{QCD} \sim 1 fm^{-1}$ , as it was assumed in 70's. Instead we have semi-classical small-size instantons and very thin QCD strings. The instanton radius peaks around the  $\rho \sim 1/3 fm$  ([2], for recent lattice data see [4] and general review [5]). String energy (action) is concentrated in a radius of .2 fm (.4 fm) in transverse directions [6]. Both are small compared to typical hadronic size, suggesting a *substructure* inside hadrons.

A snapshot of parton distribution in a transverse plane inside the nucleon should look like indicated in Fig.(1), for different  $x$  regions. These parton clusters originate from “scars” in the vacuum, being perturbed by external objects – valence quarks and strings, and therefore they must have the same transverse dimensions. One expects that these images of constituent quarks, diquarks and strings should be best seen at some intermediate  $x$ , before hadrons become black disks at very small  $x$  (high energies). The non-perturbatively produced sea quarks supposed to be more concentrated inside the constituent quarks<sup>1</sup>, while the string is supposed to be gluonic. A diquark cluster is also believed to be an instanton effect [5]. Finally, strong  $\langle \bar{q}q \rangle$  modification inside the nucleon should result in additional small density of sea quarks and gluons filling the whole disk (shown by light grey in Fig.1).

If one prefers to use the language of hadrons rather than fields, existence of two distinct components can be viewed as being due to two different scales for glueball and pion clouds, respectively. However, using hadronic description in transverse plane (or  $t$ -channel) is probably not very useful, because such complicated and coherent field configurations as instantons and strings can hardly be discussed well in this way.

These qualitative ideas were discussed in literature for long time. Constituent quarks as clusters were discussed e.g. in [2], and “scars” of two strings is behind Nachtmann-Dosch model [10] of high energy hadron-hadron scattering. But they are still mostly ignored by high energy practitioners, who think about partons as being randomly distributed inside the hadronic disk.

How can we tell whether such substructure really exists experimentally? The simplest process we have is deep-inelastic scattering (DIS). In the target frame, it can be viewed as a scattering of a dipole-like  $\bar{q}q$  objects with variable size  $d \approx 2/Q$ , where  $Q$  is the momentum transfer<sup>2</sup>. The small- $d$  dipoles measure

<sup>1</sup>This concentration should be enhanced for the polarized part of the sea. Sea quarks are found to be polarized *opposite* to valence quark (and the nucleon), as the instanton-based mechanism demands [7].

<sup>2</sup>This estimate works better for longitudinally polarized virtual photons, while for trans-

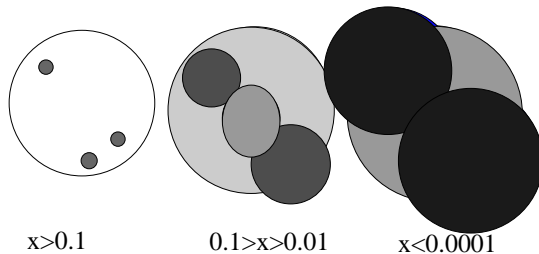


Figure 1: A snapshot of the parton density in a nucleon at different  $x$ , in an impact parameter plane.

only the average gluonic field in the target<sup>3</sup>, but dipoles with  $d$  comparable to substructure scales  $\rho$  indicated above should show a nontrivial behavior of their cross section  $\sigma(d)$  on  $d$ . pQCD predicts that  $\sigma(d)/d^2$  is constant at small  $d$ , but above some critical value  $d > d_c \approx \rho$  we expect this ratio to drop, because such large dipoles start to miss the “black spots” in the target.

Phenomenologically, when HERA data were translated into such dipole cross section [12, 13] similar behavior  $\sigma(d)$  was indeed found, and at the right scale for  $d$ . In principle however, two different effects can be responsible for it. One, emphasized in [12, 13] especially for very small  $x$  is *saturation* of the cross section, when the unitarity bound (or “blackness”) is reached. Our idea suggest similar behavior of  $\sigma(d)$ , but even at larger  $x \sim 10^{-2}$  where the spots are still “grey” and the dipole cross section is not large enough to need shadowing corrections. In fact, the particular parameterization used in [13] have exactly this feature for *all*  $x$ . Further work is needed here to understand real magnitude of both effects, due to shadowing and substructure, respectively.

Another source of information about parton correlations in transverse plane is diffraction. Clearly an inhomogeneous distribution we advocate enhances it, as compared to the homogeneous disk: the blackness of spots in higher, and there are more edges at which diffraction may take place. Studies of how diffraction depends on  $Q^2$  at HERA is therefore very important. One possible parameterization is hard-plus-soft pomerons [14], which places the boundary between soft and hard contributions at  $Q^2 \sim 4 - 10 \text{ GeV}^2$ , or  $d \sim d_c$  we speak about.

More generally, like it happened for form-factors at similar  $Q^2$  some time ago<sup>4</sup>, after new round of works we may be forced to reassess where exactly DIS

verse ones large  $d$  tail is more significant, see [13].

<sup>3</sup>For exact definition see [11].

<sup>4</sup> Although  $Q^2$  dependence of form-factors roughly follow perturbative power counting rules, their absolute magnitude significantly exceeds the pQCD predictions. Non-perturbative

is truly perturbative. The previous paradigm – the leading twist dominance down to  $Q$  as low as  $Q^2 \sim 0.5 \text{ GeV}^2$  – seems now oversimplified.

Hadron-hadron collisions are of course much more complicated. Studies of inelastic diffraction in  $\pi p$ ,  $pp$ ,  $pA$  and collision of two nuclei [15] definitely show very large  $O(1)$  fluctuations of the total nucleon cross section. It is clear that it is completely incompatible with the picture of a grey parton disk filled with multiple independent partons, with dozens of degrees of freedom involved. Only *very few* degrees of freedom may drive those fluctuations. One obvious suspect is the distance between constituent quarks, or the length of the string in Fig.1. Another is intermittent blackness of each constituent quark, presumably related to “twinkling” character of the field strength distribution in the instanton vacuum.

The pomeron itself is known to be a small-size object in transverse plane, as can be inferred from the fact that most of the  $t$ -dependence of  $pp$  scattering is explained by nucleon form-factors, and also from  $\alpha' \sim (2 \text{ GeV})^{-2}$ . It means when we see diffractive scattering of two nucleons, those are related with diffraction of their small parts<sup>5</sup>.

The history of pomeron goes back 40 years: it still works very well[8], new discoveries are being made (such as existence of a pomeron polarization vector [9]), but its microscopic dynamics still lacks theoretical understanding. pQCD promised to explain the “hard pomeron” [16], and although recent calculation [17] of the next-to-leading correction may put it in doubt, hopefully there will be a way out<sup>6</sup>. Non-perturbative approaches include incarnations of the old multiperipheral model (e.g. [18]): but they neither provide clear cut explanations of why particular hadrons should be used, nor give convincing quantitative predictions. They also have no connection with Law-Nussinov 2-gluon exchange model, which seems to be a very natural starting point explaining constant (not growing with  $s$ ) part of the cross section.

New model for growing cross section and soft pomeron have been recently proposed by Kharzeev and Levin (KL) [1]. It includes (i) a ladder made of two  $t$ -channel longitudinal gluons, as in perturbative approach; (ii) while the  $s$ -channel “rungs” of the ladder being replaced by production of *scalar* physical states. Their main motivation was to apply some known non-perturbative matrix elements, such as gluon-to-  $\pi\pi$  transition near threshold. Furthermore,

approaches, such as instanton-based calculation of the pion form-factor [22], are in quantitative agreement with data.

<sup>5</sup>Except at extremely very high energies, when the whole disk becomes black.

<sup>6</sup>Similar things happened before: let me just recall an optimistic example which may be unfamiliar to many high energy physicists. Free energy for high temperature QCD has negative  $O(\alpha_s)$  correction of modest magnitude: but higher order ones show much larger corrections of different sign. However (unlike the pomeron intercept) one can calculate this free energy non-perturbatively on the lattice. The result is about 15% below free gas, close to the  $O(\alpha_s)$  term. All high order corrections apparently canceled out! Furthermore, there are indications that some resummed or improved perturbation theory exists, in which those cancellations happen explicitly.

they inserted “realistic” scalar spectral density instead of perturbative one below some mass  $M_0$ , suggested a schematic model for it and estimated a resulting value for the pomeron intercept<sup>7</sup>  $\Delta \approx 0.08$ , close to the phenomenological value [8]. Their input was however very limited and therefore they have treated gluonic scalar spectral density in a very schematic way. It explains correctly qualitative features of the result (e.g. its dependence on the number of colors and flavors  $N_c, N_f$ ), but one may question the accuracy of the resulting numbers. Below we (i) provide further motivation for the KL approach, and also (ii) improve on their crude schematic model in several respects. We explain why one should modify the expression (1) for the pomeron intercept, and get its meaningful chiral limit. We use other information about scalar spectral density, with different glueball parameters. We found that quark-related (pion) contribution to the pomeron intercept  $\Delta$  is reduced, and the glueball one is enhanced compared to KL numbers, making them comparable at the end.

2. To motivate the approach, let us start with the following question: How using hadrons in a ladder diagram can be consistent with the statements made above, that the pomeron exchanges take place between *small parts* of the colliding hadrons, such as constituent quarks or strings?

Well, it depends: different hadrons have different sizes! For example, multiple papers use chiral Lagrangians and pions propagating *inside* the nucleon<sup>8</sup>. It follows from chiral Lagrangian that it can be used for momenta up to the so called chiral scale  $p < \Lambda_\chi \approx 1 \text{ GeV}$ . Note: it is the small pion size which matters here, not its mass. We will have similar situation for glueballs below.

With this in mind, we can move to the next question: Why is it reasonable to single out scalar  $O^{++}$  gg channel? (Apart of the fact that we know few related coupling constants.) Because both its prominent resonances – the  $\bar{q}q$  state  $\sigma$  and the scalar glueball we call  $G_0$  – are very small-size.  $\sigma$  is the pion’s brother and its interaction is covered by the same chiral scale. Remarkably,  $G_0$  is even smaller, with  $R_{G_0} \approx .2 fm$  [19, 21]. It is the smallest hadron we know, setting a record of its kind. It should be possible to construct some effective  $G_0$  Lagrangian, applicable below some momentum scale  $M_0$ . (For first attempts to build it see [23], for discussion of the magnitude of  $M_0$  see below.) The qualitative reason why it is so compact is very strong instanton-induced attraction in the scalar channel due to small-size instantons. As shown in [2, 21], in this case diluteness of the instanton vacuum  $(\rho/R)^4 \sim (1/3)^4$  (where  $R \approx 1 fm$  is instanton mean separation) is compensated by classical enhancement factor  $(8\pi^2/g^2(\rho))^2 \sim 10^2$ .

Other channels do not have this feature. For example, another gg channel one may think of is tensor  $2^{++}$ . However instanton field do not have such

<sup>7</sup>  $\Delta$  enters the total hadronic cross section energy dependence as  $\sigma(s) \sim s^\Delta$ .

<sup>8</sup> Well known ultimate model of that kind was suggested by Skyrme: in it a nucleon is made out of pions entirely. As noted by Witten, large  $N_c$  makes the nucleon static and pion field classical. Nevertheless, the  $R_N \Lambda_\chi \gg 1$  condition is still needed to justify the model, and because this parameter is not really large, Skyrme model cannot be very accurate at any  $N_c$ .

component, and so they do not act in it. Consequently the tensor glueball has normal hadronic size,  $R_{2++} \approx .8fm$  according to [19], and it would be meaningless to consider its propagation inside a nucleon. As experience with quark vector channels (where the situation is similar) shows, in such cases we have multi-hadron spectral density dual to perturbative one down to low scale  $M_{2++} \ll M_0$ .

KL have shown that the scalar channel contribute the following “non-perturbative part” to the pomeron intercept:

$$\Delta = \frac{18\pi^2}{b^2} \int \frac{dM^2}{M^6} (\rho_{phys}(M^2) - \rho^{pert}(M^2)) \quad (1)$$

where  $\rho_{phys}(M^2)$ ,  $\rho^{pert}(M^2)$  are physical and perturbative (gg cut) spectral densities respectively. The integrand in (1) is non-zero only for  $M < M_0$  because at  $M > M_0$  pQCD works and two spectral densities become identical.

3.Let us discuss properties of scalar spectral density  $\rho_{phys}(M^2)$ , first in gluodynamics and then in QCD.

We use the same normalization of the probing operator as KL

$$\theta_\mu^\mu = \frac{\beta(g)}{2g} F^{a\alpha\beta} F_{\alpha\beta}^a \simeq -\frac{bg^2}{32\pi^2} F^{a\alpha\beta} F_{\alpha\beta}^a; \quad (2)$$

so that its perturbative spectral density is

$$\rho_{pert}(M) = \frac{1}{4096} \frac{b^2 g^4 (N_c^2 - 1) M^4}{\pi^6} \quad (3)$$

The spectral density should obey the low energy theorem [3]

$$\int \frac{dM^2}{M^2} [\rho_{phys}(M^2) - \rho_{pert}(M^2)] = -4 \langle 0 | \theta_\mu^\mu(0) | 0 \rangle \quad (4)$$

As alsomphasized in [1], this relation has historically provided the first indication that there should be large non-perturbative scale in scalar gg channel. Later more direct and quantitative assessment of this effect was proposed [2], based on small-size instantons.

In gluodynamics  $\rho(M)$  is dominated by the contribution of the scalar glueball  $G_0$ , the lightest (and therefore stable) particle of this theory. Although phenomenologically its assignment to the observed scalar resonances is confused by mixing with  $\bar{q}q$  resonances and is still under debate, both multiple lattice works and the instanton model [21] point toward  $M_{G_0} = 1.5 - 1.7 GeV$ . As mentioned already, even more important is its small size, which leads to a remarkably large coupling constant to gg current<sup>9</sup>, which according to [21] is

$$\lambda_0 = \langle 0 | g^2 G_{\mu\nu}^2 | G_0 \rangle \approx 16. \pm 2 \text{ GeV}^3 \quad (5)$$

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<sup>9</sup>We remind the reader that the units in the gluodynamics is traditionally defined by setting the string tension to be the same as in QCD.

Substituting glueball contribution to the spectral density  $\rho_{G_0} = (b/(32\pi^2))^2 \lambda_0^2 \delta(M^2 - M_{G_0}^2)$  one finds the following contribution to the pomeron intercept

$$\Delta_{G_0} \approx .03 \quad (6)$$

This is not yet the complete non-perturbative part: one still has to subtract the “missing” perturbative contribution for  $M < M_0$ .

As in [3], we determine  $M_0$  from a “duality” sum rule<sup>10</sup>

$$\int^{M_0} (\rho_{G_0} - \rho_{pert}) \frac{dM^2}{M^4} = 0 \quad (7)$$

which ensures that correlator at small distances is not changed by changing from  $\rho_{pert}$  to  $\rho_{phys}$ . Solving it for  $M_0$ , one gets<sup>11</sup>  $M_0 \approx 4/, GeV$ . For a boundary between hadronic and partonic descriptions it is unexpectedly high scale indeed.

The resulting contribution of “missing perturbative states” below  $M_0$  leads to negative contribution to  $\Delta$  of about -.01. In total, we got  $\Delta_{gluodynamics} \approx .02$ , about twice the value estimated by KL.

4. Now we return to the real world with light quarks, and consider the sigma (or  $\pi\pi$ ) contribution. The major input of the KL paper is the  $\pi\pi$  coupling at small M. It follows from the scale anomaly for chiral Lagrangian [24]

$$\rho_{\pi\pi}^{M \rightarrow 0} = \frac{3M^4}{32\pi^2} \quad (8)$$

which is larger than  $\rho_{pert}$  because there is no  $g^2$  [1]. The KL schematic model assumed that  $\rho(M) = \rho_{\pi\pi}(M)$  for *all*  $M < M_0$  (see the dashed line in Fig.(2)). But this assumption cannot be true for larger masses, because the pions are known to interact strongly in this channel, forming the famous scalar  $\sigma$  resonance.

Since we know its parameters, low energy  $\pi\pi$  contribution into the correlation function in question can be easily reconstructed. Consider the isovector vector ( $\rho$ -meson) channel, for which the spectral density is well known from  $e + e -$  collisions and  $\tau$  lepton decay. In this case the  $\pi\pi$  contribution at the threshold is trivial: the pion coupling is just the pion charge. Due to pion attractive interaction, the spectral density grows from threshold, till it reaches the peak - the  $\rho$ -meson. Its magnitude can therefore be fixed from the “vector dominance”, using the normalization to the M=0 point and known  $\rho$  meson width.

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<sup>10</sup>Rather than from the low energy theorem (4). With the simple sharp cutoff we use one cannot satisfy both, so we select duality because it closer to the integral we ultimately need. However the correlator calculated in [21] is in exact agreement with this theorem.

<sup>11</sup>In KL paper significantly smaller number  $M_0 \approx 2.2 GeV$  was used for QCD. It is dual to only  $\sigma$  meson contribution, without a glueball.

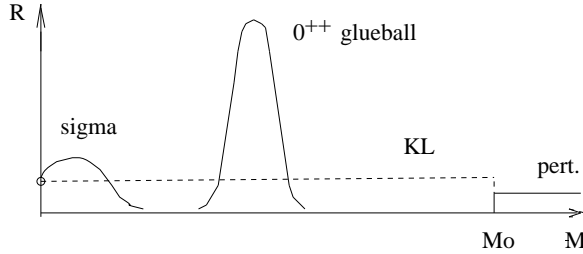


Figure 2: Schematic representation of the spectral density, as  $R = \rho(M)/\rho_{pert}(M)$  versus the total mass  $M$ . The small circle at  $M=0$  is the chiral Lagrangian prediction, the dashed line is the KL model, the solid line is our spectral density, which includes the contribution of the glueball, sigma meson and hadrons dual to perturbative gluons at large  $M$ .

Similar “sigma-dominance” should work even better, because this resonance is very wide  $m_\sigma \sim \Gamma_\sigma$ .

$$\rho_\sigma = \frac{3\pi^2 M^4}{32} \frac{M_\sigma^4}{(M^2 - M_\sigma^2)^2 + M^2 \Gamma_\sigma^2} \quad (9)$$

Naive integration from the threshold ( $2m_\pi$ ) gives large contribution to pomeron intercept  $\Delta_\sigma^{naive} \approx .09$ , where we have used  $M_\sigma = .6 \text{ GeV}$ ,  $\Gamma_\sigma = .4 \text{ GeV}$ . (Coincidentally it is close to what KS got in their paper without sigma resonance, integrating till (their)  $M_0$ .)

However, this large contribution cannot be correct, because it rests heavily on smallness of  $m_\pi$ . In the chiral limit, when quarks and pions become massless, there is no threshold at  $2m_\pi$ , and the KL  $\Delta$  simply diverges.

To resolve this problem, we should know the distribution over the so called *intrinsic* parton transverse momenta<sup>12</sup>. If gluonic partons are indeed nothing else but expansion of classical field of an instanton, their transverse momenta are related to basic non-perturbative scale  $\rho$ , the mean size of QCD instantons. If so, one gets its right magnitude  $\langle p_t^2 \rangle \approx (0.6 \text{ GeV})^2$ , and also predicts strong cut-offs, *both* at high and low momenta. The former is because the instanton is a finite-size object, leading to a form-factor  $\sim \exp(p_t \rho)$ . The latter happens because the instanton field  $A_\mu^a \sim \eta_{\mu\nu}^a x_\nu$  has a vortex-like shape with changing sign, so that its projection to constant (or long-wavelength) field vanishes<sup>13</sup>.

<sup>12</sup>Deduced e.g. from  $p_t$  distribution of Drell-Yan pairs, after perturbative effects due to extra gluon emission at large dilepton mass  $M$  are subtracted.

<sup>13</sup>In fact, there are experimental indications from diffractive dissociation cross section for transverse virtual photon is dominated by large  $p_t$ , which probably implies that indeed  $xg(x, k_t) \sim k_t^2$  at small  $k_t$ , see details in [20].



Therefore, gg collisions with small invariant mass are suppressed.

Referring to quantitative calculation elsewhere [25], here we simply modify the KL expression of (1) with a logarithmic accuracy, introducing into the KL integral over transverse momentum  $k_t^2$  of the exchanged gluons,  $\int dk_t^2/(k_t^2 + M^2)^2$ , a finite cutoff  $k_t^{min}$ . Now one of the factors  $1/M^2$  in (1) is modified, and

$$\Delta = \frac{18\pi^2}{b^2} \int \frac{dM^2}{M^4} \frac{(\rho_{phys}(M^2) - \rho^{pert}(M^2))}{(M^2 + (k_t^{min})^2)} \quad (10)$$

The unphysical divergence in the chiral limit is no longer there. Although sensitivity to the cutoff value is formally logarithmic, it still has significant effect on the intercept. With this modification, and  $(k_t^{min})^2$  for the *pair* to be equal to single parton  $\langle p_t^2 \rangle$  mentioned, we find the  $\sigma$  (or two-pion) contribution to be

$$\Delta_\sigma|_{realistic} \approx .03 \quad (11)$$

It is now comparable to the glueball contribution discussed above. This outcome is in fact more natural than the KS numbers, 0.08 and 0.01 for  $\sigma, G_0$  effects, because they are  $O(N_f^2/N_c^2)$  (for  $N_f = 2$ ) and  $O(1)$ , respectively.

**5.Summary and discussion.** Finally, combining all contributions together we get our final estimate of the non-perturbative contribution to pomeron intercept, resulting from  $O^{++}$  hadronic ladder with  $M < M_0 \approx 4 \text{ GeV}$  to be

$$\Delta = .05 \pm 0.015 \quad (12)$$

now with our guessed uncertainties.

How meaningful is this result? One still has to add to it the pQCD contribution, which is [16, 17] with the region  $M < M_0$  in the scalar channel subtracted. With current uncertainties, we do not know its value, and some readers may be disappointed at this point. However, the progress is not zero. First of all, the contribution we discuss in non-perturbative  $O(\alpha_s^0)$ , enhanced by classical instanton effects  $O(\alpha_s^{-1})$  compared to perturbative result. Second, it is sufficiently close to the phenomenological value  $\Delta = 0.08$ . Third, it can be experimentally tested by itself.

Let us briefly indicate how it can be done. The discussed model claims that the ladders made of scalar resonances explains most of the growing part of the cross section, namely  $\delta\sigma_N N \approx \sigma_{NN}(s_0) * 0.05 \log(s/s_0)$ . It means quantitative predictions about multiplicities of these scalar resonances  $\sigma, G_0$ , both in inelastic collision and their double diffractive production. Sigmas are wide and distorted by low  $p_t$  cut we do not undersnad well: so they are difficult to trace down. The glueball however shows up as relatively narrow resonance  $f_0(1500)$ , found in the double diffractive production. It has good experimental signature: large  $\eta\eta$  and  $\eta\eta'$  decay modes. So its growing production with energy may even affect s-dependence of the  $\eta/\pi$  ratio. But the hottest thing to understand is the

azimuthal distribution of nucleons in double diffractive production: in fact the WA102 data (discussed in details in [26]) show surprisingly different distribution for “glueball-type” resonances (such as  $f_0(1500)$  we discuss) from that for quark-based mesons.

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